

Mirror left-right symmetry

Pei-Hong Gu*

Max-Planck-Institut für Kernphysik, Saupfercheckweg 1, 69117 Heidelberg, Germany

We propose a novel $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ left-right symmetric model where the standard model fermion and Higgs fields are $SU(2)_L$ doublets or $SU(2)$ singlets while their mirror partners are $SU(2)_R$ doublets or $SU(2)$ singlets. The scalar fields also include a real singlet for dark matter and two $SU(2)$ triplets for seesaw. The mixing between the standard model and mirror fermions is forbidden by a $Z_2 \times Z'_2$ discrete symmetry. The mirror charged fermions can decay into their standard model partners with the dark-matter scalar while the mirror neutrinos can decay into the mirror charged fermions through the right-handed gauge interactions. Our model can have new implications on the strong CP problem, leptogenesis, collider phenomenology and dark matter detection.

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The $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ left-right symmetric theory [1], motivated by restoring the parity invariance, has become one of the most attractive proposals beyond the $SU(3)_c \times SU(2)_L \times U(1)_Y$ standard model (SM). In the existent left-right symmetric models, the SM left-handed fermions are placed in the $SU(2)_L$ doublets as they are in the SM while the SM right-handed fermions plus the right-handed neutrinos are placed in the $SU(2)_R$ doublets. After the left-right symmetry is spontaneously broken down to the electroweak symmetry, we can obtain the lepton number violation for the seesaw [2] mechanism. In the seesaw context [2, 3], the leptogenesis [4] for generating the baryon asymmetry in the universe can be realized by the Yukawa and scalar interactions [4, 5]. The left-right symmetric model for the universal seesaw scenario [6], where additional $SU(2)$ -singlet fermions with heavy masses are introduced to construct the Yukawa couplings with the $SU(2)$ -doublet fermions and Higgs, can provide a solution to the strong CP problem if the parity symmetry is imposed [7].

In this paper we shall realize the left-right symmetry in a new way where the SM fermion and Higgs fields are the $SU(2)_L$ doublets or $SU(2)$ singlets while their mirror [8] partners are the $SU(2)_R$ doublets or $SU(2)$ singlets. As a consequence of the parity symmetry, we can know the ratio among the mirror charged fermion masses from the SM charged fermion mass spectrum. Furthermore, the mirror right-handed neutrinos and the SM left-handed neutrinos can have the Majorana mass matrices with a same texture through the type-II seesaw [3]. The parity can also guarantee a vanishing strong CP phase from the θ -vacuum of the QCD and the mass matrices of the SM and mirror quarks. The mirror charged fermions can directly decay into their SM partners with a dark-matter scalar while the mirror neutrinos can decay into the lighter mirror charged fermions by exchanging the right-handed charged gauge boson. So, we can have a leptogenesis with the CP asymmetry formulated by the SM neutrino masses, lepton and quark mixing matrices. For a left-right symmetry breaking at $\mathcal{O}(10^8 \text{ GeV})$, the

first-generation mirror charged fermions can be verified by the ongoing and future colliders. In this case the dark-matter scalar can be found at the colliders even if it is not sensitive to the dark matter direct detection experiments.

The field content is summarized in Table I. Under the parity symmetry, the fields transform as

$$\begin{aligned} q_L &\leftrightarrow q'_R, \quad d_R \leftrightarrow d'_L, \quad u_R \leftrightarrow u'_L, \quad l_L \leftrightarrow l'_R, \quad e_R \leftrightarrow e'_L, \\ \xi_L &\leftrightarrow \xi_R, \quad \phi_L \leftrightarrow \phi_R, \quad \chi \leftrightarrow \chi, \\ G_\mu^a &\leftrightarrow G_\mu^a, \quad W_{L_\mu}^a \leftrightarrow W_{R_\mu}^a, \quad B_\mu \leftrightarrow B_\mu. \end{aligned} \quad (1)$$

The kinetic terms are

$$\begin{aligned} \mathcal{L}_K = & -\frac{1}{4}G_{\mu\nu}^a G^{a\mu\nu} - \frac{1}{4}W_{L(R)\mu\nu}^a W_{L(R)}^{a\mu\nu} - \frac{1}{4}B_{\mu\nu}B^{\mu\nu} \\ & + \text{Tr}\{[D_\mu \xi_{L(R)}]^\dagger D^\mu \xi_{L(R)}\} + [D_\mu \phi_{L(R)}]^\dagger D^\mu \phi_{L(R)} \\ & + \frac{1}{2}\partial_\mu \chi \partial^\mu \chi + i\bar{q}_L \not{D} q_L + i\bar{q}'_R \not{D} q'_R + i\bar{d}_R \not{D} d_R \\ & + i\bar{d}'_L \not{D} d'_L + i\bar{u}_R \not{D} u_R + i\bar{u}'_L \not{D} u'_L + i\bar{l}_L \not{D} l_L \\ & + i\bar{l}'_R \not{D} l'_R + i\bar{e}_R \not{D} e_R + i\bar{e}'_L \not{D} e'_L, \end{aligned} \quad (2)$$

where the $SU(2)$ doublets and singlets have the covariant derivatives,

$$D_\mu = \partial_\mu + ig_1 \frac{B-L}{2} B_\mu + ig_2 \frac{\tau_a}{2} W_{L_\mu(R_\mu)}^a + ig_3 \frac{\lambda_a}{2} G_\mu^a, \quad (3)$$

while the $SU(2)$ -triplet Higgs scalars have

$$D_\mu \xi_{L(R)} = (\partial_\mu + ig_1 B_\mu) \xi_{L(R)} + ig_2 W_{L_\mu(R_\mu)}^a \left[\frac{\tau_a}{2}, \xi_{L(R)}\right]. \quad (4)$$

The full scalar potential is

$$\begin{aligned} V = & \mu_{\xi_{L(R)}}^2 \text{Tr}[\xi_{L(R)}^\dagger \xi_{L(R)}] + \lambda_\xi \{\text{Tr}[\xi_{L(R)}^\dagger \xi_{L(R)}]\}^2 \\ & + \lambda'_\xi \text{Tr}[\xi_{L(R)}^\dagger \xi_{L(R)}] \text{Tr}[\xi_{L(R)} \xi_{L(R)}] \\ & + \lambda''_\xi \text{Tr}(\xi_L^\dagger \xi_L) \text{Tr}(\xi_R^\dagger \xi_R) + \mu_{\phi_{L(R)}}^2 \phi_{L(R)}^\dagger \phi_{L(R)} \\ & + \lambda_\phi [\phi_{L(R)}^\dagger \phi_{L(R)}]^2 + 2\lambda'_\phi \phi_{L(R)}^\dagger \phi_{L(R)} \phi_R^\dagger \phi_R + \frac{1}{2}\mu_\chi^2 \chi^2 \\ & + \frac{1}{4}\lambda_\chi \chi^4 + 2\lambda_{\phi\xi} \phi_{L(R)}^\dagger \phi_{L(R)} \text{Tr}[\xi_{L(R)}^\dagger \xi_{L(R)}] \\ & + 2\lambda'_{\phi\xi} \phi_{L(R)}^\dagger \phi_{L(R)} \text{Tr}[\xi_{R(L)}^\dagger \xi_{R(L)}] \\ & + \rho_{\xi_{L(R)}} [\phi_{L(R)}^T i\tau_2 \xi_{L(R)} \phi_{L(R)} + \text{H.c.}] \\ & + \lambda_{\chi\xi} \chi^2 \text{Tr}[\xi_{L(R)}^\dagger \xi_{L(R)}] + \lambda_{\chi\phi} \chi^2 \phi_{L(R)}^\dagger \phi_{L(R)} \end{aligned} \quad (5)$$

Field	$SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$	$Z_2 \times Z'_2$
$q_L = \begin{bmatrix} u_L \\ d_L \end{bmatrix}$	$(\mathbf{3}, \mathbf{2}, \mathbf{1}, +\frac{1}{3})$	$(-, +)$
$q'_R = \begin{bmatrix} u'_R \\ d'_R \end{bmatrix}$	$(\mathbf{3}, \mathbf{1}, \mathbf{2}, +\frac{1}{3})$	$(+, -)$
d_R	$(\mathbf{3}, \mathbf{1}, \mathbf{1}, -\frac{2}{3})$	$(-, +)$
d'_L	$(\mathbf{3}, \mathbf{1}, \mathbf{1}, -\frac{2}{3})$	$(+, -)$
u_R	$(\mathbf{3}, \mathbf{1}, \mathbf{1}, +\frac{4}{3})$	$(-, +)$
u'_L	$(\mathbf{3}, \mathbf{1}, \mathbf{1}, +\frac{4}{3})$	$(+, -)$
$l_L = \begin{bmatrix} \nu_L \\ e_L \end{bmatrix}$	$(\mathbf{1}, \mathbf{2}, \mathbf{1}, -1)$	$(-, +)$
$l'_R = \begin{bmatrix} \nu'_R \\ e'_R \end{bmatrix}$	$(\mathbf{1}, \mathbf{1}, \mathbf{2}, -1)$	$(+, -)$
e_R	$(\mathbf{1}, \mathbf{1}, \mathbf{1}, -2)$	$(-, +)$
e'_L	$(\mathbf{1}, \mathbf{1}, \mathbf{1}, -2)$	$(+, -)$
$\phi_L = \begin{bmatrix} \phi_L^0 \\ \phi_L^- \end{bmatrix}$	$(\mathbf{1}, \mathbf{2}, \mathbf{1}, -1)$	$(+, +)$
$\phi_R = \begin{bmatrix} \phi_R^0 \\ \phi_R^- \end{bmatrix}$	$(\mathbf{1}, \mathbf{1}, \mathbf{2}, -1)$	$(+, +)$
$\xi_L = \begin{bmatrix} \frac{1}{\sqrt{2}}\xi_L^+ & \xi_L^{++} \\ \xi_L^0 & -\frac{1}{\sqrt{2}}\xi_L^+ \end{bmatrix}$	$(\mathbf{1}, \mathbf{3}, \mathbf{1}, +2)$	$(+, +)$
$\xi_R = \begin{bmatrix} \frac{1}{\sqrt{2}}\xi_R^+ & \xi_R^{++} \\ \xi_R^0 & -\frac{1}{\sqrt{2}}\xi_R^+ \end{bmatrix}$	$(\mathbf{1}, \mathbf{1}, \mathbf{3}, +2)$	$(+, +)$
χ	$(\mathbf{1}, \mathbf{1}, \mathbf{1}, 0)$	$(-, -)$
G_μ^a	$(\mathbf{8}, \mathbf{1}, \mathbf{1}, 0)$	$(+, +)$
$W_{L\mu}^a$	$(\mathbf{1}, \mathbf{3}, \mathbf{1}, 0)$	$(+, +)$
$W_{R\mu}^a$	$(\mathbf{1}, \mathbf{1}, \mathbf{3}, 0)$	$(+, +)$
B_μ	$(\mathbf{1}, \mathbf{1}, \mathbf{1}, 0)$	$(+, +)$

TABLE I: Field content.

without any complex parameters. The allowed Yukawa interactions only include

$$\begin{aligned} \mathcal{L}_Y = & -y_d(\bar{q}_L \tilde{\phi}_L d_R + \bar{q}'_R \tilde{\phi}_R d'_L) - y_u(\bar{q}_L \phi_L u_R + \bar{q}'_R \phi_R u'_L) \\ & -y_e(\bar{l}_L \tilde{\phi}_L e_R + \bar{l}'_R \tilde{\phi}_R e'_L) - \chi(h_d \bar{d}_R d'_L + h_u \bar{u}_R u'_L \\ & + h_e \bar{e}_R e'_L) - \frac{1}{2}f(\bar{l}'_L i\tau_2 \xi_L l_L + \bar{l}'_R i\tau_2 \xi_R l'_R) + \text{H.c.} \end{aligned} \quad (6)$$

with h_d , h_u and h_e being hermitian and f being symmetric. We should keep in mind that the mass terms of the $SU(2)$ -singlet fermions are forbidden by the $Z_2 \times Z'_2$ discrete symmetry.

The symmetry breaking pattern should be $SU(2)_L \times SU(2)_R \times U(1)_{B-L} \xrightarrow{\langle \phi_R \rangle} SU(2)_L \times U(1)_Y \xrightarrow{\langle \phi_L \rangle} U(1)_{em}$,

where the vacuum expectation values (VEVs) are

$$\begin{aligned} \langle \phi_{R(L)} \rangle &= \langle \phi_{R(L)}^0 \rangle = \frac{1}{\sqrt{2}} v_{R(L)} \\ \text{with } v_{R(L)}^2 &\simeq -\frac{\lambda_\phi \mu_\phi^2 v_{R(L)}^2 - \lambda'_\phi \mu_\phi^2 v_{L(R)}^2}{\lambda_\phi^2 - \lambda_\phi'^2}. \end{aligned} \quad (7)$$

The $SU(2)$ -triplet Higgs scalars can pick up the smaller VEVs,

$$\begin{aligned} \langle \xi_{R(L)} \rangle &= \langle \xi_{R(L)}^0 \rangle = u_{R(L)} \simeq -\frac{\rho_{\xi_{R(L)}} v_{R(L)}^2}{2M_{\xi_{R(L)}}^2} \ll v_{R(L)} \\ \text{with } M_{\xi_{R(L)}}^2 &= \mu_{\xi_{R(L)}}^2 + \lambda_{\phi\xi} v_{R(L)}^2 + \lambda'_{\phi\xi} v_{L(R)}^2. \end{aligned} \quad (8)$$

In general, the four VEVs are complex and contain four phases. However, two phases can be eliminated by the $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge symmetry while the others can be set zero by minimizing the scalar potential. So, the four VEVs are all real. Note we have to require the mass terms $\mu_{\phi_R}^2 \neq \mu_{\phi_L}^2$, which can arise from a spontaneous parity violation [9], for softly breaking the parity. Otherwise, we will obtain the unaccepted $v_R = v_L$. The softly broken parity can also allow $\mu_{\xi_R}^2 \neq \mu_{\xi_L}^2$ and $\rho_{\xi_R} \neq \rho_{\xi_L}$. For the hierarchical $u_{R(L)} \ll v_{R(L)}$, most fractions of $\text{Im}(\phi_{R,L}^0)$ and $\phi_{R,L}^\pm$ while tiny fractions of $\text{Im}(\xi_{R,L}^0)$ and $\xi_{R,L}^\pm$ become the longitudinal components of the massive gauge bosons $W_{R,L}^\pm$, Z' and Z , which will be clarified later. So, we can simply take the $SU(2)$ -doublet Higgs scalars to be

$$\phi_{R(L)} = \frac{1}{\sqrt{2}} \begin{bmatrix} v_{R(L)} + h_{R(L)} \\ 0 \end{bmatrix}, \quad (9)$$

with the mass terms,

$$V \supset \lambda_\phi v_{R(L)}^2 h_{R(L)}^2 + 2\lambda'_\phi v_L v_R h_L h_R. \quad (10)$$

As for the real scalar χ , it should have a positive mass term to respect the unbroken $Z_2 \times Z'_2$ discrete symmetry,

$$m_\chi^2 = \mu_\chi^2 + \lambda_{\chi\phi} [v_{R(L)}^2] + 2\lambda_{\chi\xi} [u_{R(L)}^2] > 0. \quad (11)$$

After the symmetry breaking, there will be two charged gauge bosons,

$$\begin{aligned} W_{R(L)}^\pm &= \frac{1}{\sqrt{2}} [W_{R(L)}^1 \mp iW_{R(L)}^2] \quad \text{with} \\ M_{W_{R(L)}}^2 &= \frac{g_2^2}{4} [v_{R(L)}^2 + 4u_{R(L)}^2], \end{aligned} \quad (12)$$

two massive neutral gauge bosons,

$$\begin{aligned} Z' &= Z_R \cos \alpha - Z_L \sin \alpha, \quad Z = Z_R \sin \alpha + Z_L \cos \alpha \\ \text{with } M_{Z',Z}^2 &= \frac{g_2^2 [v_{R(L)}^2 + 8u_{R(L)}^2] \cos^2 \theta_W}{8 \cos 2\theta_W} \left\{ 1 \right. \\ &\quad \left. \pm \sqrt{1 - \frac{4(v_L^2 + 8u_L^2)(v_R^2 + 8u_R^2) \cos 2\theta_W}{[v_{R(L)}^2 + 8u_{R(L)}^2]^2 \cos^4 \theta_W}} \right\} \\ \Rightarrow M_Z^2 &\simeq \frac{g_2^2 (v_L^2 + 8u_L^2)}{4 \cos^2 \theta_W}, \quad M_{Z'}^2 \simeq \frac{g_2^2 (v_R^2 + 8u_R^2) \cos^2 \theta_W}{4 \cos 2\theta_W} \\ &\quad \text{for } v_L^2 + 8u_L^2 \ll v_R^2 + 8u_R^2, \end{aligned} \quad (13)$$

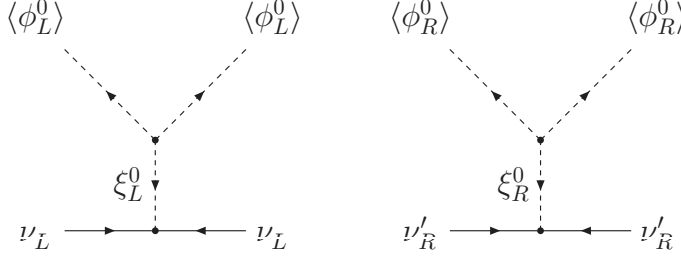


FIG. 1: The left- and right-handed type-II seesaw.

and a massless photon,

$$A = B\sqrt{\cos 2\theta_W} + W_L^3 \sin \theta_W + W_R^3 \sin \theta_W. \quad (14)$$

Here we have defined [10]

$$\sin^2 \theta_W = \frac{g_1^2}{2g_1^2 + g_2^2}, \quad (15)$$

$$Z_R = -B \tan \theta_W + W_R^3 \sec \theta_W \sqrt{\cos 2\theta_W}, \quad (16a)$$

$$Z_L = -B \tan \theta_W \sqrt{\cos 2\theta_W} + W_L^3 \cos \theta_W - W_R^3 \sin \theta_W \tan \theta_W, \quad (16b)$$

and

$$\tan 2\alpha = \frac{2 \sin^2 \theta_W \sqrt{\cos 2\theta_W}}{\frac{v_L^2 + 8u_L^2}{v_L^2 + 8u_L^2} \cos^4 \theta_W + \sin^4 \theta_W - \cos 2\theta_W}. \quad (17)$$

The SM and mirror charged fermions can obtain the Dirac masses through their Yukawa couplings with the $SU(2)$ -doublet Higgs scalars,

$$m_f = \frac{1}{\sqrt{2}} y_f v_L, \quad M_{f'} = \frac{1}{\sqrt{2}} y_{f'} v_R, \quad (18)$$

which yields

$$\begin{aligned} \frac{v_R}{v_L} &= \frac{M_{d'}}{m_d} = \frac{M_{s'}}{m_s} = \frac{M_{b'}}{m_b} = \frac{M_{u'}}{m_u} = \frac{M_{c'}}{m_c} = \frac{M_{t'}}{m_t} \\ &= \frac{M_{e'}}{m_e} = \frac{M_{\mu'}}{m_\mu} = \frac{M_{\tau'}}{m_\tau}. \end{aligned} \quad (19)$$

As for the SM left-handed neutrinos and the mirror right-handed neutrinos, they can obtain the Majorana masses through the left- and right-handed type-II seesaw [3], respectively,

$$m_\nu = f u_L, \quad M_{\nu'} = f u_R. \quad (20)$$

The relevant diagram is shown in Fig. 1. The SM and mirror neutrino mass eigenvalues should obey

$$\frac{u_R}{u_L} = \frac{M_{\nu'_1}}{m_{\nu_1}} = \frac{M_{\nu'_2}}{m_{\nu_2}} = \frac{M_{\nu'_3}}{m_{\nu_3}}. \quad (21)$$

Note that the unitary quark mixing matrix, i.e. the Cabibbo-Kobayashi-Maskawa [11] (CKM) matrix V and

the unitary lepton mixing matrix, i.e. the Maki-Nakagawa-Sakata [12] (MNS) matrix U in the SM sector should be identified with those in the mirror sector. Since the $Z_2 \times Z'_2$ symmetry forbids the mixing between the SM and mirror fermions, we will not have the one-loop $W_L - W_R$ mixing [13], which always exists in the conventional left-right symmetric models. Instead, the SM and mirror quarks associated with the real scalar can mediate a two-loop $W_L - W_R$ mixing as shown in Fig. 2. We can estimate the two-loop mass term to be

$$\begin{aligned} \delta M_{W_L W_R}^2 &\simeq \frac{3g_2^2 |V_{tb}|^2 h_{u_{33}} h_{d_{33}} M_{t'} M_{b'} m_t m_b}{2(16\pi^2)^2 \Lambda^2 (= M_{t'}^2)} \\ &\simeq \frac{3g_2^2 |V_{tb}|^2 h_{u_{33}} h_{d_{33}} m_b^2}{2(16\pi^2)^2} \ll m_b^2, \end{aligned} \quad (22)$$

which means a negligible $W_L - W_R$ mixing.

The non-perturbative QCD Lagrangian now should be

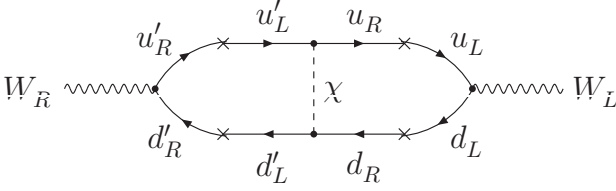
$$\mathcal{L}_{QCD} \supset \bar{\theta} \frac{g_3^2}{32\pi^2} G\tilde{G} \quad \text{with} \quad \bar{\theta} = \theta + \text{ArgDet}(M_u M_d) \quad (23)$$

where θ is the original QCD phase while M_u and M_d are the mass matrices of the SM and mirror down- and up-type quarks, respectively,

$$\begin{aligned} \mathcal{L} &\supset -[\bar{d}_L, \bar{d}'_L] M_d \begin{bmatrix} d_R \\ d'_R \end{bmatrix} - [\bar{u}_L, \bar{u}'_L] M_u \begin{bmatrix} u_R \\ u'_R \end{bmatrix} + \text{H.c.} \\ \text{with } M_d &= \begin{bmatrix} \frac{y_d v_L}{\sqrt{2}} & 0 \\ 0 & \frac{y'_d v_R}{\sqrt{2}} \end{bmatrix}, \quad M_u = \begin{bmatrix} \frac{y_u v_L}{\sqrt{2}} & 0 \\ 0 & \frac{y'_u v_R}{\sqrt{2}} \end{bmatrix}. \end{aligned} \quad (24)$$

Clearly, the θ -term should be removed as a result of the parity invariance. Furthermore, the real determinants $\text{Det}(M_d)$ and $\text{Det}(M_u)$ should induce a zero $\text{ArgDet}(M_u M_d)$. We hence can obtain a vanishing strong CP phase $\bar{\theta}$ [14]. In the absence of the mixing between the SM and mirror fermions, the strong CP phase will keep zero at loop level, unlike that in the universal seesaw scenario [7].

As shown in Fig. 3, the right-handed charged gauge boson can mediate the three-body decays of the mirror neutrinos into the mirror charged fermions as long as the kinematics is allowed. For example, if the mirror neutrinos are only heavier than the first-generation mirror

FIG. 2: The two-loop $W_L - W_R$ mixing.

charged fermions, their decay width should be

$$\begin{aligned} \Gamma_{\nu'_i} &= \Gamma_{\nu'_i \rightarrow e'_R + u'_R + d'_R} + \Gamma_{\nu'_i \rightarrow e'^c_R + u'^c_R + d'_R} \\ &\simeq \frac{g_2^4 |V_{ud}|^2 |U_{ei}|^2}{2^{10} \pi^3} \frac{M_{\nu'_i}^5}{M_{W_R}^4} \text{ for } M_{\nu'_i}^2 \ll M_{W_R}^2, \end{aligned} \quad (25)$$

In the presence of the nonzero CP phases in the unitary MNS matrix, the mirror neutrino decays can generate a lepton asymmetry in the mirror charged leptons if the decaying mirror neutrino is not heavier than the mirror tau lepton. For demonstration, we consider the case where at least one of the three mirror neutrinos can only decay into the first-generation mirror charged fermions. The CP asymmetry can be calculated at two-loop order,

$$\begin{aligned} \varepsilon_{\nu'_i} &= \frac{\Gamma_{\nu'_i \rightarrow e'_R + u'_R + d'_R} - \Gamma_{\nu'_i \rightarrow e'^c_R + u'^c_R + d'_R}}{\Gamma_{\nu'_i \rightarrow e'_R + u'_R + d'_R} + \Gamma_{\nu'_i \rightarrow e'^c_R + u'^c_R + d'_R}} \\ &= \begin{cases} \frac{g_2^4 |V_{ud}|^2}{48 \pi^3} \sum_{j \neq i} \frac{\text{Im}[(U_{ei}^* U_{ej})^2]}{|U_{ei}|^2} \frac{m_{\nu_i}}{m_{\nu_j}} & \text{for } M_{\nu'_i}^2 \ll M_{\nu'_j}^2 \ll M_{W_R}^2, \\ \frac{g_2^4 |V_{ud}|^2}{128 \pi^3} \sum_{j \neq i} \frac{\text{Im}[(U_{ei}^* U_{ej})^2]}{|U_{ei}|^2} \frac{m_{\nu_i} m_{\nu_j}}{m_{\nu_j}^2 - m_{\nu_i}^2} & \text{for } M_{\nu'_i}^2 \simeq M_{\nu'_j}^2 \ll M_{W_R}^2 \text{ and } |M_{\nu'_j}^2 - M_{\nu'_i}^2|^2 \gg M_{\nu'_i} M_{\nu'_j} \Gamma_{\nu'_i} \Gamma_{\nu'_j}. \end{cases} \end{aligned} \quad (26)$$

Since we have known [15, 16]

$$\begin{aligned} V_{ud} &= 0.97428 \pm 0.00015, \quad s_{12}^2 = 0.312^{+0.017}_{-0.015}, \\ s_{13}^2 &= \begin{cases} 0.010^{+0.009}_{-0.006} & \text{for normal hierarchy,} \\ 0.013^{+0.009}_{-0.007} & \text{for inverted hierarchy,} \end{cases} \\ \frac{m_{\nu_2}^2 - m_{\nu_1}^2}{10^{-3} \text{ eV}^2} &= \begin{cases} 2.45 \pm 0.09 & \text{for normal hierarchy,} \\ -2.34^{+0.10}_{-0.09} & \text{for inverted hierarchy,} \end{cases} \\ \frac{m_{\nu_2}^2 - m_{\nu_1}^2}{10^{-5} \text{ eV}^2} &= 7.59^{+0.20}_{-0.18}, \quad g_2^4 = 32 G_F^2 M_{W_L}^4 = 0.182, \\ U_{e1} &= c_{12} c_{13} e^{i \frac{\alpha_1}{2}}, \quad U_{e2} = s_{12} c_{13} e^{i \frac{\alpha_2}{2}}, \quad U_{e3} = s_{13} e^{-i \delta}, \end{aligned} \quad (27)$$

the CP asymmetry will only depend on four parameters: one of the neutrino mass eigenvalues $m_{\nu_{1,2,3}}$, the Dirac CP phase δ and the Majorana CP phases $\alpha_{1,2}$. For example, we can input the best fit values to read

$$\varepsilon_{\nu'_1} = 4.11 \times 10^{-7} [\sin(\alpha_2 - \alpha_1) - 0.0057 \sin(\alpha_1 + 2\delta)]$$

for the normal hierarchy with $m_{\nu_1} = 10^{-4} \text{ eV}$, (28a)

$$\varepsilon_{\nu'_3} = 1.63 \times 10^{-7} [\sin(\alpha_1 + 2\delta) + 0.45 \sin(\alpha_2 + 2\delta)]$$

for the inverted hierarchy with $m_{\nu_3} = 10^{-4} \text{ eV}$. (28b)

If the left-right symmetry is broken at $\mathcal{O}(10^8 \text{ GeV})$, the first-generation mirror charged fermions can obtain the masses from the weak scale to the TeV scale. For example, according to [15]

$$\begin{aligned} m_e &= 0.511 \text{ MeV}, \quad m_u = 1.7 - 3.1 \text{ MeV}, \\ m_d &= 4.1 - 5.7 \text{ MeV}, \quad v_L = 246 \text{ GeV}, \end{aligned} \quad (29)$$

we can expect

$$\begin{aligned} M_{e'} &= 208 \text{ GeV}, \quad M_{u'} = 0.69 - 1.3 \text{ TeV}, \\ M_{d'} &= 1.7 - 2.3 \text{ TeV} \quad \text{for } v_R = 10^8 \text{ GeV}. \end{aligned} \quad (30)$$

So, the first-generation mirror charged fermions can be produced at the ongoing and future colliders due to their couplings with the gauge bosons and the real scalar. Subsequently, they will decay into their SM partners with the real scalar being a missing energy,

$$f' \rightarrow f + \chi. \quad (31)$$

At the same time, we can take

$$\begin{aligned} m_{\nu_{1(3)}} &= 10^{-4} \text{ eV}, \quad M_{\nu_{1(3)}} = 10 \text{ TeV} \quad \text{for} \\ u_L &= 0.1 \text{ eV}, \quad u_R = 10^7 \text{ GeV}, \end{aligned} \quad (32)$$

to generate a desired lepton asymmetry in the mirror electron. As the mirror electron decays into the SM charged leptons with the real scalar, the mirror lepton asymmetry can be transferred to the SM sector. The SM lepton asymmetry eventually can result in a baryon asymmetry through the sphaleron [17] processes.

A stable SM-singlet scalar can annihilate into the light SM species through the s -channel exchange of the SM Higgs boson. The annihilation cross section can arrive at about 1 pb to give a right dark-matter relic density if the dark-matter scalar is at the weak scale. Meanwhile, the dark-matter scalar can scatter off the nucleons by the t -channel exchange of the SM Higgs boson. The dark-matter-nucleon scattering can be verified by the forthcoming dark matter direct detection experimental results. There have been a lot of works studying such dark-matter scalar [18]. In the present model, the real scalar χ definitely is a dark-matter scalar. However, the mirror charged fermions rather than the SM Higgs boson can dominate the dark-matter annihilation and scattering. In particular, when the mirror electron dominates the dark-matter annihilation, we can detect the dark-matter scalar at the colliders through the distinguishable decays of the mirror charged fermions even if the dark-matter-nucleon scattering is not significant. Furthermore, the mirror charged fermions and the dark-matter scalar can mediate other interesting processes such as lepton flavor violation and meson-antimeson mixing. We will study the mirror charged fermions and the dark-matter scalar in details elsewhere.

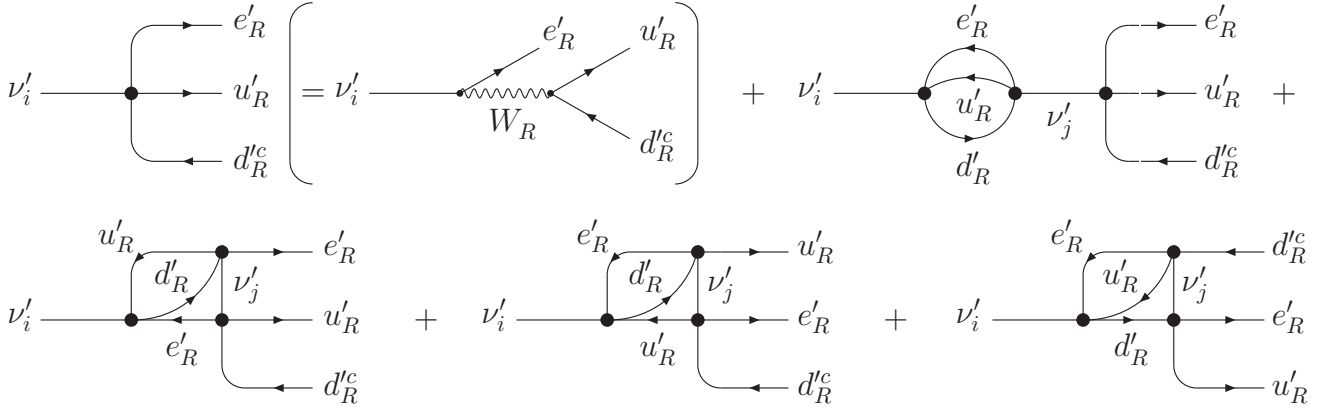


FIG. 3: The mirror Majorana neutrinos decay into the mirror charged fermions at two-loop order. The CP conjugation is not shown for simplicity.

To summarize, we have proposed a mirror left-right symmetric model with a softly broken parity symmetry. Benefited from the parity symmetry, the ratio among the mirror charged fermion masses can equal that among the SM charged fermion masses while the Majorana mass matrices of the mirror neutrinos and the SM neutrinos can have a same texture. Our model can solve the strong CP problem without the axion, like the left-right symmetric model for the universal seesaw. The mirror charged fermions can decay into their SM partners with the dark-matter scalar while the mirror neutrinos can decay into the mirror charged fermions through the right-handed gauge interactions. We thus can realize a novel leptogenesis scenario where the CP asymmetry is fully determined by the neutrino mass matrix. Although the left-right symmetry cannot be broken at the accessible TeV scale, it is possible to examine the mass spectrum of the first-generation mirror charged fermions at the colliders for a left-right symmetry breaking scale of the order of $\mathcal{O}(10^8 \text{ GeV})$. Our dark-matter scalar can be found at the colliders even if it has no significant signal in the dark matter direct detection experiments.

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* Electronic address: peihong.gu@mpi-hd.mpg.de

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